

Charge and colour breaking in the Constrained MSSM

Steven Abel^a and Toby Falk^b

^aTheory Division, Cern 1211, Geneva 23, Switzerland

^bDepartment of Physics, University of Wisconsin,
Madison, WI 53706, USA

October 14, 1998

Abstract

We show that the physical minimum of the Constrained MSSM is only free from dangerous charge and colour breaking minima in the region of parameter space bounded by $110 \lesssim m_{1/2} \lesssim 400$ GeV and $80 \lesssim m_0 \lesssim 170$ GeV. In the remaining regions the cosmology is severely constrained.

CERN-TH/98-322

MADPH-98-1083

hep-ph/9810297

1 Introduction

Unphysical charge and colour breaking (CCB) vacua [1, 2, 3, 4, 5, 6, 7, 8] have come under renewed scrutiny recently¹. For a number of models it has been found that CCB vacua are present in the whole of the parameter space which has not already been excluded by experiment. This is true for models where supersymmetry breaking is driven by the dilaton [2], for M -theory in which supersymmetry breaking is driven by bulk moduli fields [7, 8] and for the MSSM at the low $\tan\beta$ fixed point [5].

In this letter we extend the analysis of Ref.[5] to a complete determination of CCB bounds in the Constrained MSSM (CMSSM). The present work is partly an update of the results of Ref.[3], and from that study and also Ref.[5] it is clear that the parameter space is bounded from different directions by CCB bounds, dark matter bounds and experimental bounds. Dark matter bounds (the requirement that neutralino dark matter be consistent with a universe that is 12 Gyr old) tend to eliminate regions where m_0 is large. On the other hand, experiment (in particular Higgs and chargino searches) is eliminating regions of low $m_{1/2}$. CCB bounds however tend to ‘favour’ regions where $m_{1/2} \lesssim m_0$ or in other words high m_0 and low $m_{1/2}$. What we shall demonstrate in this letter is that there is only a small amount of remaining parameter space in the CMSSM which does *not* have global CCB minima.

Before continuing, we emphasise that there are various schools of thought regarding CCB minima. This is because the tunneling rate from the physical vacuum into any global CCB minima is extremely small, so that the physical vacuum is essentially stable on the lifetime of the observable universe. The authors of Ref.[3], for example, invoke a principle that ‘the cosmological constant is zero in the *global* minimum’ in order to explain the vanishing cosmological constant, thus requiring that the standard model minimum be deeper than any CCB minima. If this were the case, then CCB minima would indeed be a severe problem and would impose severe constraints on the MSSM. By contrast, some other authors take a minimalist approach; if there is no chance of tunneling, then CCB minima are not a problem. However, one has to explain how the physical vacuum was chosen over the much wider CCB minima [4, 6]. In fact this does naturally occur in some models, but not in others (*e.g.* supergravity models which possess a Heisenberg symmetry [9], including no-scale models of supergravity [10]). As the vacuum choice depends on unknown details of our cosmological history (*e.g.* the inflationary potential) we think that CCB minima should ultimately be regarded as a constraint on early cosmology rather than particle physics. We also emphasize that since tunneling between vacua is so slow as to be irrelevant, it is more appropriate to flag models which contain any CCB minima, regardless of whether they are global or local.

We begin by describing the most dangerous directions, assuming the usual R -parity invariant superpotential of the MSSM,

$$W_{MSSM} = h_U Q H_2 U^c + h_D Q H_1 D^c + h_E L H_1 E^c + \mu H_1 H_2, \quad (1)$$

¹In Refs.[5, 7] the bounds which will be of most interest here were, in deference to historical precedent, referred to as Unbounded From Below (UFB) bounds. This is a confusing misnomer since the directions are not, in most cases, unbounded from below, and we revert to calling them CCB bounds.

and a degenerate pattern of supersymmetry breaking with universal scalar mass² parameters (m_0^2), trilinear couplings (A_0) and gaugino masses ($m_{1/2}$) at the GUT or Planck scale. The notation is the same as in Ref.[5].

As emphasised in Ref.[6], the flat directions in the MSSM are a direct result of the adoption of R -parity, to prevent the proton decaying. The dangerous F and D flat directions which will be of interest in this letter are constructed from gauge invariants involving H_2 [11, 12]. This is because its mass squared parameter, m_2^2 , appears in the potential along these directions, and it must be negative in order to drive electroweak symmetry breaking. The first example of this kind in the literature is (see Komatsu in Ref.[1])

$$L_i Q_3 D_3 ; H_2 L_i \quad (2)$$

where the suffices on matter superfields are generation indices. With the following choice of VEVs;

$$\begin{aligned} h_2^0 &= -a^2 \mu / h_{D_{33}} \\ \tilde{d}_{L_3} = \tilde{d}_{R_3} &= a \mu / h_{D_{33}} \\ \tilde{\nu}_i &= a \sqrt{1 + a^2} \mu / h_{D_{33}}, \end{aligned} \quad (3)$$

the potential along this direction is F and D -flat, and depends only on the soft supersymmetry breaking terms;

$$V = \frac{\mu^2}{h_{D_{33}}^2} a^2 (a^2 (m_2^2 + m_{L_{ii}}^2) + m_{L_{ii}}^2 + m_{d_{33}}^2 + m_{Q_{33}}^2). \quad (4)$$

At large values of $a \gg 1$ the potential is governed by the first term. Because m_2^2 is required to turn negative during the renormalisation group running (for successful electroweak symmetry breaking) the potential can develop a charge and colour breaking minimum at a scale of $few \times \mu / h_{D_j}$. Ensuring that this does not happen leads to the constraint in which we are interested.

The above is not quite, but is very close to, the deepest ‘fully optimised’ direction [1, 3]. To get the fully optimized condition we parameterize the VEV of $\tilde{\nu}_i$ as [1]

$$\tilde{\nu}_i = \gamma_L a^2 \mu / h_{D_{33}}. \quad (5)$$

Minimisation of V with respect to γ_L then gives

$$\gamma_L^2 = \frac{1 + a^2}{a^2} - \frac{2m_{L_i}^2 h_b^2}{\hat{g}^2 a^4 \mu^2} \quad (6)$$

where $\hat{g}^2 = (g'^2 + g_2^2)/2$, and a potential of

$$V = \frac{\mu^2}{h_{D_{33}}^2} a^2 (a^2 (m_2^2 + m_{L_{ii}}^2) + m_{L_{ii}}^2 + m_{d_{33}}^2 + m_{Q_{33}}^2) - \frac{m_{L_i}^4}{\hat{g}^2}. \quad (7)$$

When $\gamma_L^2 < 0$ then one should set $\gamma_L^2 = 0$ to get the most stringent condition although, for the regions of parameter space of interest here, this will never be the case. The potential typically

has a depth of $\gtrsim 10^6 m_W^4$ at the minimum. In addition, as shown in Ref.[6], the depth of the minimum is actually extremely sensitive to the choice of soft supersymmetry breaking parameters, so that, for most reasonable choices of parameters, the difference between Eq.(4) and Eq.(7) is negligible. In Ref.[6], emphasis was also placed on the closeness of the usual condition (that the physical vacuum be the global minimum) to the more relevant condition (that the physical vacuum be the only minimum). There is generally an extremely thin, but cosmologically interesting, region of parameter space between the ‘allowed’ and ‘disallowed’ regions, where there is a CCB minimum which is *local*.

In order to obtain the bound we now need to take account of the renormalisation group running of the mass-squared parameters between the weak and GUT scales. To do this we shall assume that the largest mass, and therefore the appropriate scale at which to evaluate the parameters is $\phi = h_{U33} \langle h_2^0 \rangle$. This minimises the top quark contributions to the effective potential at one-loop. Further corrections to the potential are assumed to be small. As shown Ref.[3], this approximation is adequate for determining CCB bounds on the supersymmetry breaking parameters (although they also note a curious bifurcation in behaviour when one varies the scale at which the parameters are evaluated).

In the above potentials, $\langle h_2^0 \rangle = -a^2 \mu / h_{D33, E33}$ so that the Eq.(4) is of the form

$$V = \frac{M_{GUT}^2}{h_{U33}^2} \hat{\phi} (\hat{\phi} A + B/b) \quad (8)$$

where $A = m_2^2(\phi) + m_{L_{ii}}^2(\phi)$, B is the other combination of mass-squared parameters (also evaluated at ϕ) which appears in the potentials above,

$$\hat{\phi} = \phi / M_{GUT} \quad (9)$$

and

$$b(\phi) = \frac{M_{GUT} h_{D33}}{h_{U33} \mu} \quad (10)$$

for the LQD , LH_2 direction described above, or

$$b(\phi) = \frac{M_{GUT} h_{E33}}{h_{U33} \mu} \quad (11)$$

for the equally dangerous LLE , LH_2 direction. The traditional (no global CCB minima) bound is saturated by $V = V' = 0$; the non-trivial solution is therefore also a solution to $\tilde{V} = \tilde{V}' = 0$ where

$$\tilde{V} = \hat{\phi} A + B/b. \quad (12)$$

Hence, only the GUT scale parameters plus the parameter b enter the traditional bound as was pointed out in Ref.[6]. In fact the bound always becomes *more* restrictive with increase in b , since this decreases the positive contribution to the potential.

2 The analysis

2.1 CCB Minima

First we shall examine the properties of the CCB minima when we vary the soft supersymmetry breaking parameters. We will not however calculate the tunneling rates which is generally a complicated numerical task. In this instance, it is made particularly difficult by the presence of more than one gauge invariant involved in the flat directions. (As an aside we note that, for the directions involving only the single UDD operator examined in Ref.[4], a very accurate analytic approximation can be found for the tunneling rate.) However, from dimensional considerations and estimates of tunneling rates in Refs.[4, 6], it is clear that the lifetime for these directions will be much longer than the age of the universe.

We can make additional observations on the stability of physical vacuum by adopting the approach of Ref.[3] in the light of Ref.[6]. In the latter it was pointed out that the CCB bounds are least restrictive when $A_0 = -m_{1/2}$ ². Hence in Fig. 1 we show contours of the VEV of the CCB minimum in the $m_0, m_{1/2}$ parameter space, for $\tan\beta = 2, 10$ and $\mu < 0$, and taking $A_0 = -m_{1/2}$ and 0. Here we have plotted $\log_{10}(\text{VEV}/246 \text{ GeV})$ for the LLE, LH_2 direction. Above the solid contour there are no global CCB minima. In the thin shaded region (approximately 10 GeV wide) above the solid contour, there are CCB minima which are only local, and the physical minimum is still the global minimum. In Ref.[7] a general survey was made of all the possible dangerous directions in field space, and this direction was found to give the severest bounds (at least for the CMSSM) so henceforth we shall only be considering this. The figure illustrates that the VEVs of CCB minima are smallest when the CCB bound is close to being saturated. This leads to the rather counter-intuitive fact that the physical vacuum is least stable in these regions. When the bound is strongly violated the VEVs become very large and the lifetime of the physical vacuum increases. The CCB bounds are least restrictive when $A \approx -m_{1/2}$, and they are close to the analytic approximation found in Ref.[6]; if we define $\tilde{m}_0 = m_0/m_{1/2}$ then the analytic approximation is given by

$$\tilde{m}_0^2 > \frac{f(3\tilde{m}_0^2) - g(3\tilde{m}_0^2)(1 - \rho_p)}{4 - 3\rho_p} \quad (13)$$

where

$$\begin{aligned} f(x) &= 1.43 - 0.16x + 0.02x^2 \\ g(x) &= 2.94 - 0.2x + 0.02x^2 \\ 1/\rho_p &= 1 + 3.17(\sin^2\beta - \sin^2\beta^{QFP}). \end{aligned} \quad (14)$$

These functions were evaluated in Ref.[7] for $m_{1/2} = 200 \text{ GeV}$ and in the one loop approximation. However we use the full two loop value of $\tan\beta^{QFP}$ which is ≈ 1.6 . (Note that the important factor here is the distance from the fixed point which is given by ρ_p . For a given value of $\tan\beta$ this obviously depends sensitively on the fixed point value, $\tan\beta^{QFP}$, so we

²In our sign conventions, the 3-4 element of the neutralino mass matrix is $-\mu$, and the mixing term in the stop mass matrix is $m_t(A_t + \mu \cot\beta)$.

cannot expect the analytic approximation to be better than $\sim 15\%$.) For $\tan\beta = 2, 10$ and $A_0 = -m_{1/2}$ we find the bounds $m_0 \gtrsim 130, 40$ respectively. The first compares favourably with the bounds at $m_{1/2} = 200$ GeV found numerically in Fig. 2 but the $\tan\beta = 10$ bound is found to be larger than the analytic estimate. This suggests that the bottom quark Yukawa and/or the two loop contributions to the beta functions are already contributing significantly to the bounds at $\tan\beta = 10$.

2.2 Experimental Constraints

We now discuss the experimental and cosmological bounds on the CMSSM parameter space. Recent runs at LEP at center-of-mass energies of 172 and 183 GeV have excluded large areas of the CMSSM parameters space, and subsequent runs at ~ 190 and 200 GeV will push the bounds even further. In the CMSSM, the dominant constraints at moderate to high $\tan\beta$ come from searches for chargino pair production, and modulo a small loophole which can occur when the mass of the sneutrino is close to the chargino mass, the experimental bounds saturate the kinematic limit of $m_{\chi^\pm} \sim 91$ GeV. Chargino iso-mass contours of 91 and 100 GeV, representing the current and projected LEP 200 chargino mass bounds, respectively, are displayed in Fig. 2 as dashed lines in the $m_{1/2} - m_0$ plane, for two representative values of $\tan\beta$ and both signs of μ . Also shown as a dotted line is the current LEP183 slepton mass bound [13], which is roughly 84 GeV for large $m_{\tilde{l}_R} - m_{\tilde{\chi}}$.

At low $\tan\beta$, the dominant CMSSM constraint comes from searches for Higgs production at LEP183. Not only are the experimental bounds strongest at low $\tan\beta$, roughly 90 GeV at $\tan\beta = 2$, but here the tree level Higgs mass is also smallest ($m_h^{\text{tree}} \approx m_Z \cos 2\beta$ for $m_A \gg m_Z$). Radiative corrections to the Higgs mass [14], which depend logarithmically on the sfermion masses, must then be very large, leading to strong lower bounds on the masses of the sfermions, and in particular the stops. However, the extraction of the radiatively corrected Higgs mass in the MSSM has an uncertainty of ~ 2 GeV, so we conservatively take $m_h > 88$ GeV as our experimental lower limit at low $\tan\beta$. The LEP183 Higgs bounds are shown in Fig. 2 as a dot-dashed line. The entire displayed region of Fig. 2a is excluded by the Higgs mass constraint, while none of the displayed region in Fig. 2d is in conflict with the current Higgs bound. Of course the Higgs mass corrections are very sensitive to $\tan\beta$ for $\tan\beta$ near the quasi-fixed point, and, additionally, the experimental limit falls for $\tan\beta > 2$, and so the Higgs constraint moves quickly to the left for $\tan\beta > 2$. The chargino bound, on the other hand, moves to the right, and for $100 \text{ GeV} < m_0 < 200 \text{ GeV}$, the two bounds together exclude $m_{1/2} \lesssim 110$ GeV for all $\tan\beta$ and both signs of μ .

2.3 Cosmological Constraints

Over most of the CMSSM parameter space, the lightest supersymmetric particle (LSP) is a bino-like lightest neutralino $\tilde{\chi}$. R -parity ensures that the LSP is stable over cosmological time scales, and in the CMSSM, the LSP typically has a cosmologically interesting relic density [15, 16]. In much of the parameter space, in fact, the relic abundance of neutralinos is so large that it is in conflict with the observed age of the universe, $t_U > 12$ Gyr, and

the corresponding upper limit of $\Omega_\chi h^2 < 0.3$ can be used to exclude large areas of m_0 and $m_{1/2}$, as follows. In the early universe, gaugino-like neutralinos annihilate predominantly via sfermion exchange. Increasing either m_0 or $m_{1/2}$ drives up the sfermion masses, lowers the annihilation rate, and increases the neutralino relic abundance. Thus an upper bound on $\Omega_\chi h^2$ translates into an upper bound on the parameters $m_0, m_{1/2}$. The experimental and cosmological constraints are nicely complementary, in part because the cosmological bounds put upper limits on the parameters for which particle searches provide lower limits.

The dark shaded area in Fig. 2 delimits the cosmologically preferred region with $0.1 < \Omega_\chi h^2 < 0.3$. The upper limit, as described above, comes from an upper bound on the age of the universe. The lower limit is more of a preference than a bound, stemming from the wish to have the neutralinos comprise a significant fraction of the dark matter. The two narrow vertical channels arise from s-channel neutralino annihilation on the Higgs and Z^0 poles (for some values of $\tan\beta$ they have merged into one large pole region). Note that the top of the shaded area intersects with the “theory” excluded area (where a stau is the LSP) at $m_{1/2} \sim 450$ GeV. Thus $\Omega_\chi h^2 < 0.3$ yields an upper bound on $m_{1/2}$ and on m_0 (except in the close vicinity of the pole region, which is largely excluded by the LEP chargino searches). At sufficiently low $\tan\beta$, the Higgs bound moves to the right of the shaded region, and the incompatibility of a Higgs lower bound of 86 GeV with the lower limit on the age of the universe excludes values of $\tan\beta$ less than 2, for $\mu < 0$, or 1.65 for $\mu > 0$ [17].

The dark solid line in Fig. 1 is reproduced in Fig. 2, where we have chosen $A_0 = -m_{1/2}$ to minimise the size of area containing CCB minima (see above and Ref.[6]). We note that there is only a restricted region in $\{m_{1/2}, m_0\}$ which is both cosmologically and experimentally viable and which is free of CCB minima, and moreover, such a region exists only for $\tan\beta > 2.3$ for $\mu < 0$ and $\tan\beta > 2.1$ for $\mu > 0$. At large $\tan\beta \gtrsim 20$, s-channel annihilation through the pseudoscalar Higgs and heavy Higgs can contribute significantly to the neutralino annihilation cross-section, and the above cosmological bounds are weakened[18]. Coannihilations with light sleptons also substantially reduces the relic abundance of neutralinos near the line $m_{\tilde{\tau}} = m_{\tilde{\chi}}$ and can provide a window at large $m_{1/2}$ [19].

3 Conclusions

In this paper we have re-examined the unphysical charge and colour breaking minima in the Constrained MSSM. In summary, we find that most of the parameter space which is not already excluded by experiment or by cosmological considerations has an unphysical vacuum which is lower than the physical one. The region of parameter space which remains is roughly

$$\begin{aligned} 110 &\lesssim m_{1/2} \lesssim 400 \text{ GeV} \\ 80 &\lesssim m_0 \lesssim 170 \text{ GeV}, \end{aligned} \tag{15}$$

although there is a narrow region at $m_{1/2} \approx 150$ (110) GeV for $\tan\beta = 10$ (3) for which higher values of m_0 are still allowed by the projected chargino searches. Future tri-lepton searches at the Tevatron should push the lower bound on $m_{1/2}$ to from 180 – 240 GeV for the regions of Fig. 2 [20] and will help close this loophole to larger m_0 .

The regions which do have an unphysical minimum, as we stated in the introduction, are not completely excluded but they must have a constrained cosmology. Some general points were made in Refs.[4, 6], so let us apply some of these considerations to this specific model. Our main requirement for any acceptable cosmological scenario is that it should explain why the physical minimum is chosen instead of the global minimum. There are currently three possible explanations;

- A high reheat temperature after inflation
- Extra contributions to the effective potential during inflation
- Breaking of R -parity

In the case that the field is driven to the origin during reheating, there is an additional constraint in all models in which supersymmetry breaking is transmitted to the physical sector gravitationally. In these models the gravitino mass is typically of the same order as the weak scale, and successful nucleosynthesis requires that

$$T_{reheat} \lesssim 10^9 \text{ GeV}. \quad (16)$$

In order to lift the unphysical minimum we should have a reheat temperature which is greater than roughly $10^{-1\pm 1} \times$ the scale of the VEV (the ‘uncertainty’ coming from the shape of the effective potential and the contribution of light particles to it). Hence a particular T_{reheat} can be effective in the region bounded by the corresponding contour in Fig. (1). From Fig. (1) we see that the condition (16) can quite easily be satisfied away from the fixed point in the CMSSM.

If the physical vacuum is chosen because of extra terms generated during inflation then there are constraints on the possible inflationary potentials. For example potentials with a Heisenberg symmetry are eliminated in this case because they are flat at tree level and the one loop corrections to the mass-squareds tend to be negative [9, 10].

If R -parity is broken explicitly [5, 6] we must choose a symmetry other than R -parity to prevent the proton from decaying, and of course lose the neutralino as a dark matter candidate (but, on the plus side, free up a large region of parameter space in Fig. (2)). Note that, as well as the lepton number violating version of R -parity violation considered in Ref.[6], the CCB minima would also be lifted by a large Majorana mass for a right handed neutrino. If R -parity is broken during a stage of pre-heating [21] then it is in principle possible to have neutralino dark matter which is stable. It is also possible that further lifting of the potential would occur in any case from extra operators required in the visible sector. Both of these mechanisms work because at large field values (*i.e.* those corresponding to the VEV of the unphysical minimum), the MSSM is no longer a good description of the effective potential, and the potential is no longer F flat along the relevant directions.

4 Acknowledgements

We would like to thank Sasha Davidson and Carlos Savoy for discussions, and the CERN theory division for hospitality.

References

- [1] J.-M. Frère, D.R.T. Jones and S. Raby, *Nucl. Phys.* **B222** (1983) 11; M. Claudson, L. Hall and I. Hinchcliffe, *Nucl. Phys.* **B228** (1983) 501; H.-P. Nilles, M. Srednicki and D. Wyler, *Phys. Lett.* **B120** (1983) 346; J-P. Derendinger and C.A. Savoy, *Nucl. Phys.* **B237** (1984) 307; H. Komatsu, *Phys. Lett.* **B215** (1988) 323; P. Langacker and N. Polonsky, *Phys. Rev.* **D50** (1994) 2199; A. Kusenko, P. Langacker and G. Segre, *Phys. Rev.* **D54** (1996) 5824; J.A. Casas and S. Dimopoulos, *Phys. Lett.* **B387** (1996) 107; J.A. Casas, [hep-ph/9707475](#) ; J.A. Casas, A. Lleyda and C. Munoz, *Nucl. Phys.* **B471** (1996) 3; J.A. Casas, A. Lleyda and C. Munoz, *Phys. Lett.* **B389** (1996) 305
- [2] J.A. Casas, A. Lleyda and C. Munoz, *Phys. Lett.* **B380** (1996) 59
- [3] H. Baer, M. Brhlik, D. Castano, *Phys. Rev.* **D54** (1996) 6944
- [4] T. Falk, K.A. Olive, L. Roszkowski and M. Srednicki, *Phys. Lett.* **B367** (1996) 183; A. Riotto and E. Roulet, *Phys. Lett.* **B377** (1996) 60; A. Strumia, *Nucl. Phys.* **B482** (1996) 24; T. Falk, K.A. Olive, L. Roszkowski, A. Singh and M. Srednicki, *Phys. Lett.* **B396** (1997) 50
- [5] S.A. Abel and B.C. Allanach, *Phys. Lett.* **B415** (371) 1997; [hep-ph/9803476](#)
- [6] S.A. Abel and C.A. Savoy, [hep-ph/9803218](#)
- [7] S.A. Abel and C.A. Savoy, [hep-ph/9809498](#)
- [8] J.A. Casas, A. Ibarra and C. Munoz, [hep-ph/9810266](#)
- [9] P. Binetruiy and M.K. Gaillard, *Phys. Lett.* **B195** (1987) 382
- [10] M.-K. Gaillard, H. Murayama, and K. Olive, *Phys. Lett.* **B355** (1995) 71
- [11] F. Buccella, J-P. Derendinger, C. A. Savoy and S. Ferrara, *Phys. Lett.* **B115** (1982) 375; R. Gatto and G. Sartori, *Phys. Lett.* **B157** (1985) 389; M. A. Luty and W. Taylor, *Phys. Rev.* **D53** (1996) 3399
- [12] T. Ghergetta, C. Kolda, S. P. Martin, *Nucl. Phys.* **B468** (1996) 37
- [13] ALEPH Collaboration, R. Barate et al., CERN preprint EP/98-077 (1998)
- [14] H.E. Haber, R. Hempfling and A.H. Hoang, *Z. Phys.* **C75** 539; see also M. Carena, M. Quiros and C.E.M. Wagner, *Nucl. Phys.* **B461** (1996) 407
- [15] H. Goldberg, *Phys. Rev. Lett.* **50** (1983) 1419
- [16] J. Ellis, J.S. Hagelin, D.V. Nanopoulos, K.A. Olive and M. Srednicki, *Nucl. Phys.* **B238** (1984) 453
- [17] J. Ellis, T. Falk, G. Ganis, K.A. Olive and M. Schmitt, *Phys. Rev.* **D58** (1998) 095002.

- [18] M. Drees and M. M. Nojiri, Phys. Rev. **D47** (1993) 376; V. Barger and C. Kao, Phys. Rev. **D57** (1998) 3131.
- [19] J. Ellis, T. Falk, and K.A. Olive, in preparation.
- [20] V. Barger and C. Kao, MAD-98-1085, in preparation.
- [21] A. Riotto, E Roulet and I Vilja, *Phys. Lett.* **B390** (1997) 73

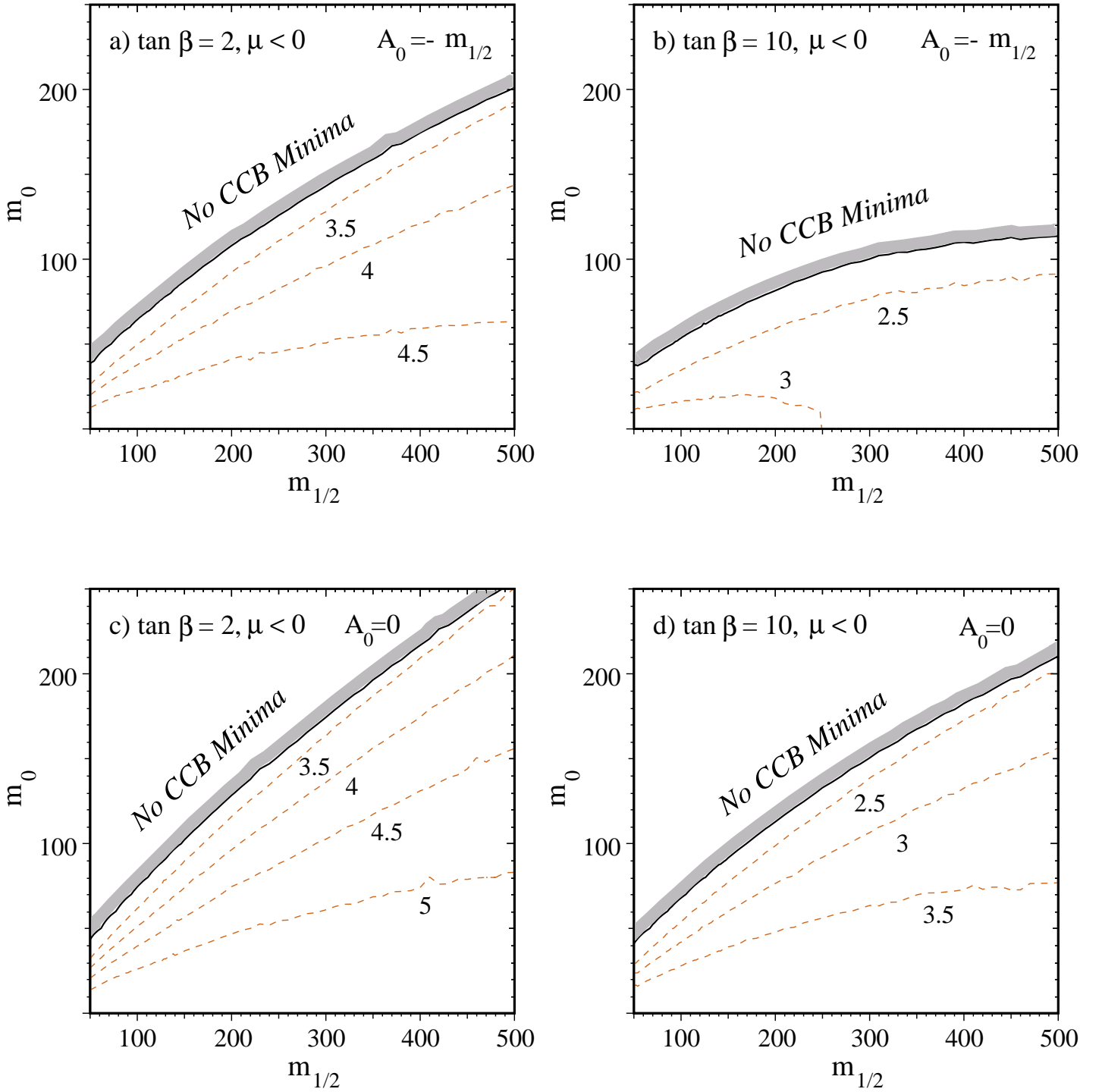


Figure 1: The location of the CCB VEV, in the direction LLE, LH_2 , for $\tan \beta = 2, 10$ and $A_0 = 0, -m_{1/2}$. The dashed lines are contours of constant $\log_{10}(\text{VEV}/246 \text{ GeV})$. In the shaded strip, the CCB minimum is not global. There are no CCB minima above the shaded region.

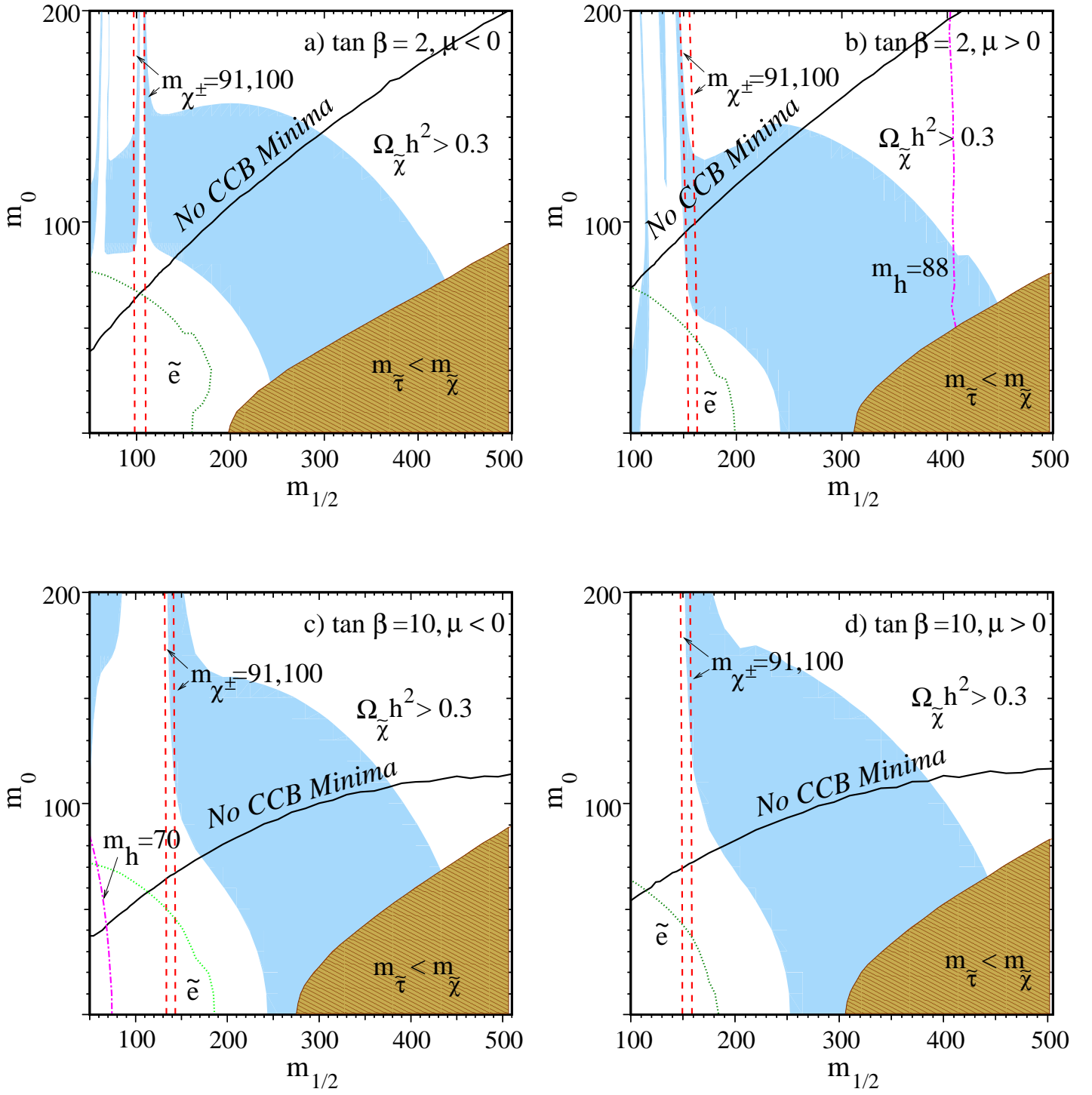


Figure 2: The combined cosmological and experimental constraints on the constrained MSSM, for $\tan \beta = 2, 10$ and both $\mu < 0$ and $\mu > 0$. The dashed contours represent current and future LEP chargino bounds, dotted contours are slepton bounds, and dot-dashed contours are Higgs bounds. The light-shaded region gives $0.1 < \Omega_{\tilde{\chi}} h^2 < 0.3$. Below the solid contour, CCB minima are present in the LLE, LH_2 direction. We have chosen $A_0 = -m_{1/2}$ to minimise the area containing CCB minima.